STRUCTURE OF THE $(0^+, 1^+)$ MESONS B_{s0} AND B_{s1}, AND THE STRONG COUPLING CONSTANTS $g_{B_{s0}BK}$ AND $g_{B_{s1}B^*K}$

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Abstract

In this article, we take the point of view that the bottomed $(0^+, 1^+)$ mesons B_{s0} and B_{s1} are the conventional $b\bar{s}$ meson, and calculate the strong coupling constants $g_{B_{s0}BK}$ and $g_{B_{s1}B^*K}$ with the light-cone QCD sum rules. The numerical values of strong coupling constants $g_{B_{s1}B^*K}$ and $g_{B_{s0}BK}$ are very large, and support the hadronic dressing mechanism. Just like the scalar mesons $f_0(980)$, $a_0(980)$, D_{s0} and axial-vector meson D_{s1}, the $(0^+, 1^+)$ bottomed mesons B_{s0} and B_{s1} may have small $b\bar{s}$ kernels of the typical $b\bar{s}$ meson size, the strong couplings to the hadronic channels (or the virtual mesons loops) may result in smaller masses than the conventional $b\bar{s}$ mesons in the potential quark models, and enrich the pure $b\bar{s}$ states with other components.

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1 Introduction

Recently, the CDF Collaboration reports the first observation of two narrow resonances consistent with the orbitally excited P-wave B_s mesons using 1 fb^{-1} of $p\bar{p}$ collisions at $\sqrt{s} = 1.96$TeV collected with the CDF II detector at the Fermilab Tevatron [1]. The masses of the two states are $M(B_{s1}) = (5829.4 \pm 0.7)$MeV and $M(B^*_{s2}) = (5839.7 \pm 0.7)$MeV, and they can be assigned as the $J^P = (1^+, 2^+)$ states in the heavy quark effective theory [2]. The D0 Collaboration reports the direct observation of the excited P-wave state B^*_{s2} in fully reconstructed decays to B^+K^-, the mass of the B^*_{s2} meson is measured to be $(5839.6 \pm 1.1 \pm 0.7)$MeV [3]. While the B_s states with spin-parity $J^P = (0^+, 1^+)$ are still lack experimental evidence.

The masses of the B_s mesons with $(0^+, 1^+)$ have been estimated with the potential quark models, heavy quark effective theory and lattice QCD [4, 5, 6, 7, 8, 9, 10, 11, 12, 13, 14, 15, 16], the values are different from each other. In our previous work [17], we study the masses of the bottomed $(0^+, 1^+)$ mesons with the QCD sum rules, and observe that the central values are below the corresponding BK and B^*K thresholds respectively. The strong decays $B_{s0} \to BK$ and $B_{s1} \to B^*K$ are kinematically forbidden, the P-wave heavy mesons B_{s0} and B_{s1} can decay through the isospin violation precesses $B_{s0} \to B_s\eta \to B_s\pi^0$ and $B_{s1} \to B^*_s\eta \to B^*_s\pi^0$, respectively [18]. The η and π^0 transition matrix is very small according to Dashen's theorem

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[19], $t_{\eta\pi} = \langle \pi^0|\mathcal{H}|\eta\rangle = -0.003\text{GeV}$, they maybe very narrow. The bottomed mesons B_{s0} and B_{s1} may have interesting feature, just like their charmed cousins D_{s0} and D_{s1}, have small $b\bar{s}$ kernels of the typical $b\bar{s}$ mesons size, strong couplings to the virtual intermediate hadronic states (or the virtual mesons loops) may result in smaller masses than the conventional $b\bar{s}$ mesons in the potential quark models, enrich the pure $b\bar{s}$ states with other components [14, 20, 21, 22, 23, 24].

In previous works, the mesons $f_0(980)$, $a_0(980)$, D_{s0} and D_{s1} are taken as the conventional $q\bar{q}$ and $c\bar{s}$ states respectively, and the values of the strong coupling constants g_{f_0KK}, g_{a_0KK}, $g_{D_{s0}DK}$ and $g_{D_{s1}D^*K}$ are calculated with the light-cone QCD sum rules [22, 23, 24, 25, 26]. The large values of the strong coupling constants support the hadronic dressing mechanism.

In this article, we take the bottomed mesons B_{s0} and B_{s1} as the conventional $b\bar{s}$ states, and calculate the values of the strong coupling constants $g_{B_{s0}BK}$ and $g_{B_{s1}B^*K}$ with the light-cone QCD sum rules, and study the possibility of the hadronic dressing mechanism in the bottomed channels.

The light-cone QCD sum rules approach carries out the operator product expansion near the light-cone $x^2 \approx 0$ instead of the short distance $x \approx 0$ while the non-perturbative matrix elements are parameterized by the light-cone distribution amplitudes (which classified according to their twists) instead of the vacuum condensates [27, 28, 29, 30, 31, 32]. The non-perturbative parameters in the light-cone distribution amplitudes are calculated by the conventional QCD sum rules and the values are universal [33, 34, 35, 36].

The article is arranged as: in Section 2, we derive the strong coupling constants $g_{B_{s0}BK}$ and $g_{B_{s1}B^*K}$ with the light-cone QCD sum rules; in Section 3, the numerical result and discussion; and in Section 4, conclusion.

2 Strong coupling constants $g_{B_{s1}B^*K}$ and $g_{B_{s0}BK}$ with light-cone QCD sum rules

In the following, we write down the definition for the strong coupling constants $g_{B_{s0}BK}$ and $g_{B_{s1}B^*K}$,

$$
\begin{aligned}
\langle B_{s1}|B^*K\rangle &= -ig_{B_{s1}B^*K}\eta^* \cdot \epsilon = -iM_A\hat{g}_{B_{s1}B^*K}\eta^* \cdot \epsilon \, , \\
\langle B_{s0}|BK\rangle &= g_{B_{s0}BK} = M_S\hat{g}_{B_{s0}BK} \, ,
\end{aligned} \tag{1}
$$

where the ϵ_α and η_α are the polarization vectors of the mesons B^* and B_{s1} respectively. The masses M_S and M_A can serve as an energy scale, we factorize the masses from the corresponding strong coupling constants $g_{B_{s0}BK}$ and $g_{B_{s1}B^*K}$ respectively.

We study the strong coupling constants with the two-point correlation functions

$\Pi_{\mu\nu}(p,q)$ and $\Pi_{\mu}(p,q)$ respectively,

$$\begin{aligned}
\Pi_{\mu\nu}(p,q) &= i\int d^4x\, e^{-iq\cdot x}\,\langle 0|T\left\{J_\mu^V(0)J_\nu^{A+}(x)\right\}|K(p)\rangle\,, \qquad (2)\\
\Pi_{\mu}(p,q) &= i\int d^4x\, e^{-iq\cdot x}\,\langle 0|T\left\{J_\mu^5(0)J^{S+}(x)\right\}|K(p)\rangle\,, \qquad (3)\\
J_\mu^V(x) &= \bar{u}(x)\gamma_\mu b(x)\,,\\
J_\mu^A(x) &= \bar{s}(x)\gamma_\mu\gamma_5 b(x)\,,\\
J_\mu^5(x) &= \bar{u}(x)\gamma_\mu\gamma_5 b(x)\,,\\
J^S(x) &= \bar{s}(x)b(x)\,, \qquad (4)
\end{aligned}$$

where the currents $J_\mu^V(x)$, $J_\mu^A(x)$, $J_\mu^5(x)$ and $J^S(x)$ interpolate the bottomed mesons B^*, B_{s1}, B and B_{s0}, respectively, the external K meson has four momentum p_μ with $p^2 = m_K^2$. The correlation functions $\Pi_{\mu\nu}(p,q)$ and $\Pi_\mu(p,q)$ can be decomposed as

$$\begin{aligned}
\Pi_{\mu\nu}(p,q) &= i\Pi_A(p,q)g_{\mu\nu} + \Pi_{A1}(p,q)(p_\mu q_\nu + p_\nu q_\mu) + \cdots\,,\\
\Pi_{\mu}(p,q) &= i\Pi_S(p,q)q_\mu + \Pi_{S1}(p,q)p_\mu + \cdots \qquad (5)
\end{aligned}$$

due to the Lorentz invariance.

According to the basic assumption of current-hadron duality in the QCD sum rules approach [33, 34, 35, 36], we can insert a complete series of intermediate states with the same quantum numbers as the current operators $J_\mu^V(x)$, $J_\mu^A(x)$, $J_\mu^5(x)$ and $J^S(x)$ into the correlation functions $\Pi_{\mu\nu}(p,q)$ and $\Pi_\mu(p,q)$ to obtain the hadronic representation. After isolating the ground state contributions from the pole terms of the mesons B^*, B_{s1}, B and B_{s0}, we get the following results,

$$\begin{aligned}
\Pi_{\mu\nu} &= \frac{\langle 0|J_\mu^V(0)\mid B^*(q+p)\rangle\langle B^*|B_{s1}K\rangle\langle B_{s1}(q)|J_\nu^{A+}(0)|0\rangle}{\left[M_{B^*}^2-(q+p)^2\right]\left[M_A^2-q^2\right]} + \cdots\\
&= -\frac{ig_{B_{s1}B^*K}f_{B^*}f_A M_{B^*}M_A}{\left[M_{B^*}^2-(q+p)^2\right]\left[M_A^2-q^2\right]}g_{\mu\nu} + \cdots\,, \qquad (6)\\
\Pi_{\mu} &= \frac{\langle 0|J_\mu^5(0)\mid B(q+p)\rangle\langle B|B_{s0}K\rangle\langle B_{s0}(q)|J^{S+}(0)|0\rangle}{\left[M_{B}^2-(q+p)^2\right]\left[M_S^2-q^2\right]} + \cdots\\
&= \frac{ig_{B_{s0}BK}f_B f_S M_S}{\left[M_{B}^2-(q+p)^2\right]\left[M_S^2-q^2\right]}(p+q)_\mu + \cdots\,, \qquad (7)
\end{aligned}$$

where the following definitions for the weak decay constants have been used,

$$\begin{aligned}
\langle 0|J_\mu^V(0)|B^*(p)\rangle &= f_{B^*}M_{B^*}\epsilon_\mu\,,\\
\langle 0|J_\mu^A(0)|B_{s1}(p)\rangle &= f_A M_A\eta_\mu\,,\\
\langle 0|J_\mu^5(0)|B(p)\rangle &= if_B p_\mu\,,\\
\langle 0|J^S(0)|B_{s0}(p)\rangle &= f_S M_S\,. \qquad (8)
\end{aligned}$$

In the following, we briefly outline the operator product expansion for the correlation functions $\Pi_{\mu\nu}(p,q)$ and $\Pi_{\mu}(p,q)$ in perturbative QCD theory. The calculations are performed at the large space-like momentum regions $(q+p)^2 \ll 0$ and $q^2 \ll 0$, which correspond to the small light-cone distance $x^2 \approx 0$ required by the validity of the operator product expansion approach. We write down the propagator of a massive quark in the external gluon field in the Fock-Schwinger gauge firstly [37],

$$\begin{aligned}\langle 0|T\{q_i(x_1)\,\bar{q}_j(x_2)\}|0\rangle &= i\int \frac{d^4k}{(2\pi)^4}e^{-ik(x_1-x_2)} \\ &\left\{\frac{\not{k}+m}{k^2-m^2}\delta_{ij} - \int\limits_0^1 dv\, g_s\, G^{\mu\nu}_{ij}(vx_1+(1-v)x_2)\right. \\ &\left.\left[\frac{1}{2}\frac{\not{k}+m}{(k^2-m^2)^2}\sigma_{\mu\nu} - \frac{1}{k^2-m^2}v(x_1-x_2)_\mu\gamma_\nu\right]\right\}\,. \end{aligned} \quad (9)$$

Substituting the above b quark propagator and the corresponding K meson light-cone distribution amplitudes into the correlation functions $\Pi_{\mu\nu}(p,q)$ and $\Pi_{\mu}(p,q)$, and completing the integrals over the variables x and k, finally we obtain the analytical results, which are given explicitly in the appendix.

In calculation, the two-particle and three-particle K meson light-cone distribution amplitudes have been used [38, 39, 40, 41], the explicit expressions are given in the appendix. The parameters in the light-cone distribution amplitudes are scale dependent and are estimated with the QCD sum rules [38, 39, 40, 41]. In this article, the energy scale μ is chosen to be $\mu = 1$GeV, to be more precise, one can choose $\mu = \sqrt{M_B^2 - m_b^2} \approx 2.4$GeV.

After straightforward calculations, we obtain the final expressions of the double Borel transformed correlation functions $\Pi_A(M_1^2, M_2^2)$ and $\Pi_S(M_1^2, M_2^2)$ at the level of quark-gluon degrees of freedom. The masses of the bottomed mesons are $M_A = 5.72$GeV, $M_S = 5.70$GeV, $M_{B^*} = 5.33$GeV and $M_B = 5.28$GeV,

$$\frac{M_A^2}{M_A^2 + M_{B^*}^2} \approx \frac{M_S^2}{M_S^2 + M_B^2} \approx 0.54\,, \quad (10)$$

there exists an overlapping working window for the two Borel parameters M_1^2 and M_2^2, it's convenient to take the value $M_1^2 = M_2^2$. We introduce the threshold parameter s_0 and make the simple replacement,

$$e^{-\frac{m_b^2+u_0(1-u_0)m_K^2}{M^2}} \rightarrow e^{-\frac{m_b^2+u_0(1-u_0)m_K^2}{M^2}} - e^{-\frac{s_0}{M^2}}$$

to subtract the contributions from the high resonances and continuum states [37], finally we obtain the sum rules for the strong coupling constants $g_{B_{s0}BK}$ and $g_{B_{s1}B^*K}$,

$$
\begin{aligned}
g_{B_{s0}BK} &= \frac{1}{f_B f_S M_S} \exp\left(\frac{M_S^2}{M_1^2}+\frac{M_B^2}{M_2^2}\right)\left\{\left[\exp\left(-\frac{\Xi}{M^2}\right)-\exp\left(-\frac{s_S^0}{M^2}\right)\right]\right. \\
& \frac{f_K m_K^2 M^2}{m_u+m_s}\left[\varphi_p(u_0)-\frac{d\varphi_\sigma(u_0)}{6du_0}\right]+\exp\left(-\frac{\Xi}{M^2}\right)\left[-m_b f_K m_K^2\int_0^{u_0} dt B(t)\right. \\
& +f_{3K} m_K^2 \int_0^{u_0} d\alpha_s \int_{u_0-\alpha_s}^{1-\alpha_s} d\alpha_g \varphi_{3K}(1-\alpha_s-\alpha_g,\alpha_g,\alpha_s)\frac{2(\alpha_s+\alpha_g-u_0)-3\alpha_g}{\alpha_g^2} \\
& -\frac{2m_b f_K m_K^4}{M^2}\int_{1-u_0}^1 d\alpha_g \frac{1-u_0}{\alpha_g^2}\int_0^{\alpha_g} d\beta \int_0^{1-\beta} d\alpha \Phi(1-\alpha-\beta,\beta,\alpha) \\
& +\frac{2m_b f_K m_K^4}{M^2}\left(\int_0^{1-u_0} d\alpha_g \int_{u_0-\alpha_g}^{u_0} d\alpha_s \int_0^{\alpha_s} d\alpha+\int_{1-u_0}^1 d\alpha_g\int_{u_0-\alpha_g}^{1-\alpha_g} d\alpha_s \int_0^{\alpha_s} d\alpha\right) \\
& \left.\left.\frac{\Phi(1-\alpha-\alpha_g,\alpha_g,\alpha)}{\alpha_g}\right]\right\}\, ,
\end{aligned}
\qquad (11)
$$

$$
\begin{aligned}
g_{B_{s1}B^*K} &= \frac{1}{f_{B^*}f_A M_{B^*}M_A}\exp\left(\frac{M_A^2}{M_1^2}+\frac{M_{B^*}^2}{M_2^2}\right)\left\{\left[\exp\left(-\frac{\Xi}{M^2}\right)-\exp\left(-\frac{s_A^0}{M^2}\right)\right]\right. \\
& f_K\left[\frac{m_b m_K^2 M^2}{m_u+m_s}\varphi_p(u_0)+\frac{m_K^2(M^2+m_b^2)}{8}\frac{d}{du_0}A(u_0)-\frac{M^4}{2}\frac{d}{du_0}\phi_K(u_0)\right] \\
& -\exp\left(-\frac{\Xi}{M^2}\right)\left[f_K m_b^2 m_K^2\int_0^{u_0}dtB(t)\right. \\
& +m_K^2\int_0^{u_0}d\alpha_s\int_{u_0-\alpha_s}^{1-\alpha_s}d\alpha_g\frac{(u_0 f_K m_K^2\Phi+f_{3K}m_b\varphi_{3K})(1-\alpha_s-\alpha_g,\alpha_s,\alpha_g)}{\alpha_g} \\
& +f_K m_K^2 M^2\frac{d}{du_0}\int_0^{u_0}d\alpha_s\int_{u_0-\alpha_s}^{1-\alpha_s}d\alpha_g\frac{(A_\parallel-V_\parallel)(1-\alpha_s-\alpha_g,\alpha_s,\alpha_g)}{2\alpha_g} \\
& -f_K m_K^2 M^2\frac{d}{du_0}\int_0^{u_0}d\alpha_s\int_{u_0-\alpha_s}^{1-\alpha_s}d\alpha_g A_\parallel(1-\alpha_s-\alpha_g,\alpha_s,\alpha_g)\frac{\alpha_s+\alpha_g-u_0}{\alpha_g^2} \\
& +f_K m_K^4\left(\int_0^{1-u_0}d\alpha_g\int_{u_0-\alpha_g}^{u_0}d\alpha_s\int_0^{\alpha_s}d\alpha+\int_{1-u_0}^{1}d\alpha_g\int_{u_0-\alpha_g}^{1-\alpha_g}d\alpha_s\int_0^{\alpha_s}d\alpha\right) \\
& \left[\frac{1}{\alpha_g}\left(3-\frac{2m_b^2}{M^2}\right)\Phi+\frac{4m_b^2}{M^2}\frac{\alpha_s+\alpha_g-u_0}{\alpha_g^2}(A_\perp+A_\parallel)\right](1-\alpha-\alpha_g,\alpha,\alpha_g) \\
& -f_K m_K^4 u_0\frac{d}{du_0}\left(\int_0^{1-u_0}d\alpha_g\int_{u_0-\alpha_g}^{u_0}d\alpha_s\int_0^{\alpha_s}d\alpha+\int_{1-u_0}^{1}d\alpha_g\int_{u_0-\alpha_g}^{1-\alpha_g}d\alpha_s\int_0^{\alpha_s}d\alpha\right) \\
& \frac{\Phi(1-\alpha-\alpha_g,\alpha,\alpha_g)}{\alpha_g} \\
& -f_K m_K^4\int_{1-u_0}^{1}d\alpha_g\int_0^{\alpha_g}d\beta\int_0^{1-\beta}d\alpha\left[\Phi(1-\alpha-\beta,\alpha,\beta)\frac{1-u_0}{\alpha_g^2}\left(4-\frac{2m_b^2}{M^2}\right)\right. \\
& \left.+\frac{4m_b^2}{M^2}\frac{(1-u_0)^2}{\alpha_g^3}(A_\parallel+A_\perp)(1-\alpha-\beta,\alpha,\beta)\right] \\
& \left.\left.+f_K m_K^4\frac{d}{du_0}\int_{1-u_0}^{1}d\alpha_g\int_0^{\alpha_g}d\beta\int_0^{1-\beta}d\alpha\Phi(1-\alpha-\beta,\alpha,\beta)\frac{u_0(1-u_0)}{\alpha_g^2}\right]\right\}, \qquad (12)
\end{aligned}
$$

where

$$
\begin{aligned}
\Xi &= m_b^2+u_0(1-u_0)m_K^2\,, \\
u_0 &= \frac{M_1^2}{M_1^2+M_2^2}\,, \\
M^2 &= \frac{M_1^2M_2^2}{M_1^2+M_2^2}\,. \qquad (13)
\end{aligned}
$$

The term proportional to the $M^4\frac{d}{du_0}\phi_K(u_0)$ in Eq.(12) depends heavily on the asymmetric coefficient $a_1(\mu)$ of the twist-2 light-cone distribution amplitude $\phi_K(u)$ in the

limit $u_0 = \frac{1}{2}$(see also the sum rules for the strong coupling constant $g_{D_{s1}D^*K}$ in Ref. [23]), if we take the value $a_1(\mu) = 0.06 \pm 0.03$ [38, 39, 40, 41], no stable sum rules can be obtained, the value of the $g_{B_{s1}B^*K}$ changes significantly with the variation of the Borel parameter M^2. In this article, we take the assumption that the u and s quarks have symmetric momentum distributions and neglect the coefficient $a_1(\mu)$.

In the heavy quark limit $m_b \to \infty$,

$$
\begin{aligned}
s_S^0 &\to m_b^2 + 2m_b\omega_S^0 \, , \\
s_A^0 &\to m_b^2 + 2m_b\omega_A^0 \, , \\
M_1^2 &\to 2m_bT_1 \, , \\
M_2^2 &\to 2m_bT_2 \, , \\
M^2 &\to 2m_bT \, , \\
M_S &\to m_b + \Lambda_1 \, , \\
M_A &\to m_b + \Lambda_1 \, , \\
M_B &\to m_b + \Lambda_0 \, , \\
M_{B^*} &\to m_b + \Lambda_0 \, , \qquad (14)
\end{aligned}
$$

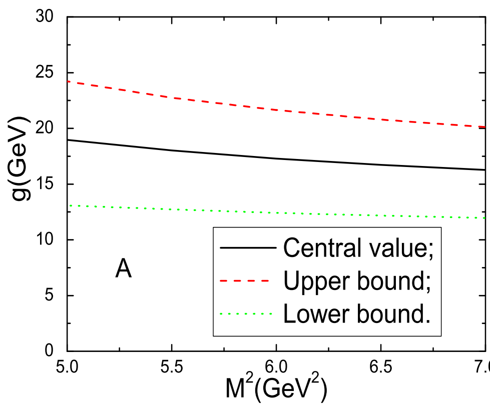
the two sum rules in Eqs.(11-12) are reduced to the following form,

$$
\begin{aligned}
g_{B_{s0}BK} &= \frac{1}{f_Bf_S}\exp\left(\frac{\Lambda_1}{T_1}+\frac{\Lambda_0}{T_2}\right)\left\{\left[1-\exp\left(-\frac{\omega_S^0}{T}\right)\right]\right. \\
&\quad \left.\frac{2f_Km_K^2T}{m_u+m_s}\left[\varphi_p(u_0)-\frac{d\varphi_\sigma(u_0)}{6du_0}\right]-f_Km_K^2\int_0^{u_0}dtB(t)\right\} \, , \qquad (15) \\
g_{B_{s1}B^*K} &= \frac{1}{f_{B^*}f_A}\exp\left(\frac{\Lambda_1}{T_1}+\frac{\Lambda_0}{T_2}\right)\left\{\left[1-\exp\left(-\frac{\omega_A^0}{T}\right)\right]\right. \\
&\quad f_K\left[\frac{2m_K^2T}{m_u+m_s}\varphi_p(u_0)+\frac{m_K^2(2T+m_b)}{8m_b}\frac{d}{du_0}A(u_0)-2T^2\frac{d}{du_0}\phi_K(u_0)\right] \\
&\quad \left.-f_Km_K^2\int_0^{u_0}dtB(t)\right\} \, , \qquad (16)
\end{aligned}
$$

where the decay constants take the behavior $f_A = \frac{C_1}{\sqrt{M_A}}$, $f_S = \frac{C_1}{\sqrt{M_S}}$, $f_B = \frac{C_2}{\sqrt{M_B}}$ and $f_{B^*} = \frac{C_2}{\sqrt{M_{B^*}}}$ (according to the definition in Eq.(8)), the C_i are some constants.

3 Numerical result and discussion

The input parameters are taken as $m_s = (140 \pm 10)$MeV, $m_u = (5.6 \pm 1.6)$MeV, $m_b = (4.7 \pm 0.1)$GeV, $\lambda_3 = 1.6 \pm 0.4$, $f_{3K} = (0.45 \pm 0.15) \times 10^{-2}\text{GeV}^2$, $\omega_3 = -1.2 \pm 0.7$, $\eta_4 = 0.6 \pm 0.2$, $\omega_4 = 0.2 \pm 0.1$, $a_2 = 0.25 \pm 0.15$ [32, 38, 39, 40, 41], $f_K = 0.160$GeV, $m_K = 0.498$GeV, $M_B = 5.279$GeV, $M_{B^*} = 5.325$GeV [42], $M_S = (5.70 \pm 0.11)$GeV, $M_A = (5.72 \pm 0.09)$GeV, $f_S = f_A = (0.24 \pm 0.02)$GeV [17],



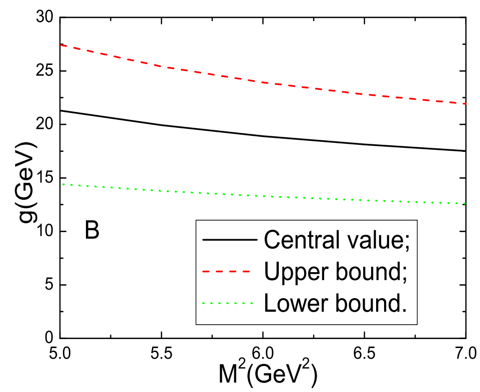


Figure 1: The strong coupling constants $g_{B_{s1}B^*K}$(A) and $g_{B_{s0}BK}$(B) with the Borel parameter M^2.

$f_{B^*} = f_B = (0.17 \pm 0.02)\mathrm{GeV}$ [32, 43, 44], $s_S^0 = (37 \pm 1)\mathrm{GeV}^2$, $s_A^0 = (38 \pm 1)\mathrm{GeV}^2$ [17], $\Lambda_0 = \frac{M_B + 3M_{B^*}}{4} - m_b = (0.6 \pm 0.1)\mathrm{GeV}$, $\Lambda_1 = \frac{M_S + 3M_A}{4} - m_b = (1.0 \pm 0.1)\mathrm{GeV}$, $\omega_S^0 = (1.6 \pm 0.1)\mathrm{GeV}$ and $\omega_A^0 = (1.6 \pm 0.1)\mathrm{GeV}$. The Borel parameters are chosen as $M^2 = (5-7)\mathrm{GeV}^2$, in this region, the values of the strong coupling constants $g_{B_{s1}B^*K}$ and $g_{B_{s0}BK}$ are rather stable, which are shown in Fig.1. In the heavy quark limit, the Borel parameters are chosen as $T = (0.7 - 1.5)\mathrm{GeV}$, in this region, the values of the strong coupling constants $g_{B_{s1}B^*K}$ and $g_{B_{s0}BK}$ are rather stable, which are shown in Fig.2.

In the limit of large Borel parameter M^2, the strong coupling constants $g_{B_{s1}B^*K}$ and $g_{B_{s0}BK}$ take up the following behaviors,

$$
\begin{aligned}
g_{B_{s0}BK} &\propto \frac{M^2 \varphi_p(u_0)}{f_B f_S}\,, \\
g_{B_{s1}B^*K} &\propto \frac{m_b M^2 \varphi_p(u_0)}{f_{B^*} f_A}\,. \qquad (17)
\end{aligned}
$$

It is not unexpected, the contributions from the two-particle twist-3 light-cone distribution amplitude $\varphi_p(u)$ are greatly enhanced by the large Borel parameter M^2, (large) uncertainties of the relevant parameters presented in above equations have significant impact on the numerical results. The contributions from the two-particle twist-2, twist-3 and twist-4 light-cone distribution amplitudes $\phi_K(u_0)$, $\varphi_\sigma(u_0)$ and $A(u_0)$ are zero due to symmetry property.

Taking into account all the uncertainties of the input parameters, finally we

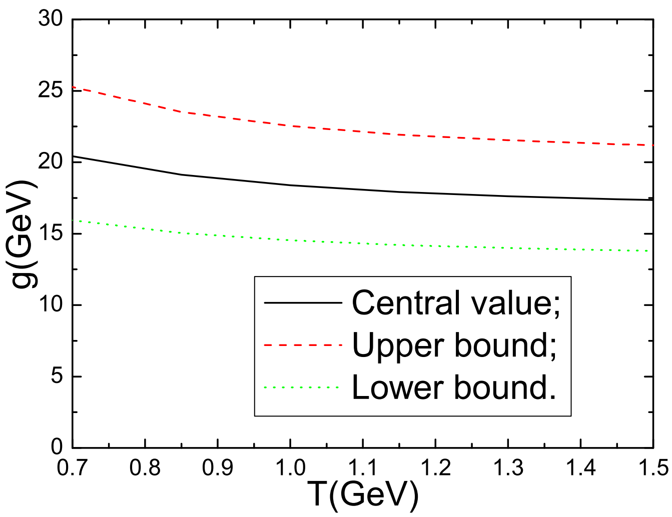


Figure 2: The strong coupling constants $g_{B_{s1}B^*K}$ and $g_{B_{s0}BK}$ with the Borel parameter T in the heavy quark limit.

obtain the numerical values of the strong coupling constants

$$\begin{aligned} g_{B_{s1}B^*K} &= (18.1 \pm 6.1)\text{GeV}\,, \\ g_{B_{s0}BK} &= (20.0 \pm 7.4)\text{GeV}\,, \\ \hat{g}_{B_{s1}B^*K} &= 3.2 \pm 1.1\,, \\ \hat{g}_{B_{s0}BK} &= 3.5 \pm 1.3 \end{aligned} \tag{18}$$

from Eqs.(11-12) and

$$g_{B_{s1}B^*K} = g_{B_{s0}BK} = (19.6 \pm 5.7)\text{GeV} \tag{19}$$

from Eqs.(15-16). The uncertainties are large, about 30%. The contributions from three-particle light-cone distribution amplitudes vanish in the heavy quark limit, the uncertainties are reduced slightly, as the dominating contributions come from the two-particle twist-3 light-cone distribution amplitude $\varphi_p(u)$.

The large values of the strong coupling constants $g_{B_{s1}B^*K}$ and $g_{B_{s0}BK}$ obviously support the hadronic dressing mechanism [45, 46, 47], the scalar meson $B_{s0}(D_{s0})$ and axial-vector meson $B_{s1}(D_{s1})$ can be taken as having small scalar and axial-vector $b\bar{s}$ ($c\bar{s}$) kernels of typical meson size with large virtual S-wave $BK(DK)$ and $B^*K(D^*K)$ cloud respectively.

In Refs.[48, 49], the authors analyze the unitarized two-meson scattering amplitudes from the heavy-light chiral lagrangian, and observe that the scalar mesons D_{s0} and B_{s0}, and axial-vector mesons D_{s1} and B_{s1} appear as the bound state poles with the strong coupling constants $g_{D_{s0}DK} = 10.203\text{GeV}$, $g_{D_{s1}D^*K} = 10.762\text{GeV}$,

	$g_{B_{s1}B^*K}$(GeV)	$g_{B_{s0}BK}$(GeV)	$g_{D_{s1}D^*K}$(GeV)	$g_{D_{s0}DK}$(GeV)
[22, 23]			10.5 ± 3.5	$9.3^{+2.7}_{-2.1}$
[48, 49]	23.572	23.442	10.762	10.203
This work	18.1 ± 6.1	20.0 ± 7.4		
This work*	19.6 ± 5.7	19.6 ± 5.7		

Table 1: Theoretical estimations of the strong coupling constants from different models, where * stands for the strong coupling constants in the heavy quark limit.

$g_{B_{s1}B^*K} = 23.572\text{GeV}$ and $g_{B_{s0}BK} = 23.442\text{GeV}$. Our numerical results for the strong coupling constants are certainly reasonable and can make robust predictions. However, we take the point of view that the mesons D_{s0}, B_{s0}, D_{s1} and B_{s1} be bound states in the sense that they appear below the corresponding DK, BK, D^*K and B^*K thresholds respectively, their constituents may be the bare $c\bar{s}$ and $b\bar{s}$ states, the virtual DK, BK, D^*K and B^*K pairs and their mixing, rather than the DK, BK, D^*K and B^*K bound states.

In Ref.[50], the authors calculate the strong coupling constants $g_{D_{s0}D_s\eta}$ and $g_{D_{s1}D_s^*\eta}$ with the light-cone QCD sum rules, then take into account $\eta - \pi^0$ mixing and calculate their pionic decay widths. The bottomed mesons B_{s0} and B_{s1} can decay through the same isospin violation mechanism, $B_{s0} \to B_s\eta \to B_s\pi^0$ and $B_{s1} \to B_s^*\eta \to B_s^*\pi^0$. We study the strong coupling constants $g_{B_{s0}B_s\eta}$ and $g_{B_{s1}B_s^*\eta}$ with the light-cone QCD sum rules and make predictions for the corresponding small decay widths [18].

4 Conclusion

In this article, we take the point of view that the bottomed mesons B_{s0} and B_{s1} are the conventional $b\bar{s}$ mesons and calculate the strong coupling constants $g_{B_{s0}BK}$ and $g_{B_{s1}B^*K}$ with the light-cone QCD sum rules. The numerical results are compatible with the existing estimations, the large values support the hadronic dressing mechanism. Just like the scalar mesons $f_0(980)$, $a_0(980)$, D_{s0} and axial-vector meson D_{s1}, the bottomed mesons B_{s0} and B_{s1} may have small $b\bar{s}$ kernels of typical $b\bar{s}$ meson size. The strong couplings to virtual intermediate hadronic states (or the virtual mesons loops) can result in smaller masses than the conventional 0^+ and 1^+ mesons in the potential quark models, enrich the pure $b\bar{s}$ states with other components.

Appendix

The analytical expressions of the $\Pi_S(p,q)$ and $\Pi_A(p,q)$ at the level of the quark-gluon degrees of freedom,

$$
\begin{aligned}
\Pi_S &= \frac{f_K m_K^2}{m_u+m_s}\int_0^1 du\frac{\varphi_p(u)}{\Delta} - m_b f_K m_K^2\int_0^1 du\int_0^u dt\frac{B(t)}{\Delta^2} \\
&+\frac{1}{6}\frac{f_K m_K^2}{m_u+m_s}\int_0^1 du\varphi_\sigma(u)\frac{d}{du}\frac{1}{\Delta} \\
&+f_{3K}m_K^2\int_0^1 dv\int_0^1 d\alpha_g\int_0^{1-\alpha_g} d\alpha_s\varphi_{3K}(\alpha_u,\alpha_g,\alpha_s)\frac{2v-3}{\Delta^2}\mid_{u=(1-v)\alpha_g+\alpha_s} \\
&-4m_b f_K m_K^4\int_0^1 dvv\int_0^1 d\alpha_g\int_0^{\alpha_g} d\beta\int_0^{1-\beta} d\alpha\frac{\Phi(1-\alpha-\beta,\beta,\alpha)}{\Delta^3}\mid_{u=1-v\alpha_g} \\
&+4m_b f_K m_K^4\int_0^1 dv\int_0^1 d\alpha_g\int_0^{1-\alpha_g} d\alpha_s\int_0^{\alpha_s} d\alpha\frac{\Phi(1-\alpha-\alpha_g,\alpha_g,\alpha)}{\Delta^3}\mid_{u=(1-v)\alpha_g+\alpha_s}\ , \\
\Pi_A &= -\frac{f_K m_b m_K^2}{m_u+m_s}\int_0^1 du\frac{\varphi_p(u)}{\Delta} + f_K m_b^2 m_K^2\int_0^1 du\int_0^u dt\frac{B(t)}{\Delta^2} \\
&+\frac{f_K}{2}\int_0^1 du\left\{\phi_K(u)\frac{d}{du}\log\Delta+\frac{A(u)m_K^2}{4}\frac{d}{du}\left[\frac{1}{\Delta}+\frac{m_b^2}{\Delta^2}\right]\right\} \\
&+m_K^2\int_0^1 dv\int_0^1 d\alpha_g\int_0^{1-\alpha_g} d\alpha_s \\
&\frac{\left[f_K m_K^2 u\Phi+f_{3K}m_b\varphi_{3K}\right](1-\alpha_s-\alpha_g,\alpha_s,\alpha_g)}{\Delta^2}\mid_{u=\alpha_s+(1-v)\alpha_g} \\
&-\frac{f_K m_K^2}{2}\int_0^1 dv\int_0^1 d\alpha_g\int_0^{1-\alpha_g} d\alpha_s\left[(1-2v)A_\parallel-V_\parallel\right](1-\alpha_s-\alpha_g,\alpha_s,\alpha_g) \\
&\frac{d}{du}\frac{1}{\Delta}\mid_{u=\alpha_s+(1-v)\alpha_g} \\
&+f_K m_K^4\int_0^1 dv\int_0^1 d\alpha_g\int_0^{1-\alpha_g} d\alpha_s\int_0^{\alpha_s} d\alpha\Phi(1-\alpha-\alpha_g,\alpha,\alpha_g) \\
&\left\{\frac{4}{\Delta^2}-\frac{4m_b^2}{\Delta^3}+u\frac{d}{du}\frac{1}{\Delta^2}\right\}\mid_{u=\alpha_s+(1-v)\alpha_g} \\
&-f_K m_K^4\int_0^1 dvv\int_0^1 d\alpha_g\int_0^{\alpha_g} d\beta\int_0^{1-\beta} d\alpha\Phi(1-\alpha-\beta,\alpha,\beta) \\
&\left\{\frac{4}{\Delta^2}-\frac{4m_b^2}{\Delta^3}+u\frac{d}{du}\frac{1}{\Delta^2}\right\}\mid_{u=1-v\alpha_g} \\
&+8f_K m_b^2 m_K^4\int_0^1 dvv\int_0^1 d\alpha_g\int_0^{1-\alpha_g} d\alpha_s\int_0^{\alpha_s} d\alpha(A_\parallel+A_\perp)(1-\alpha-\alpha_g,\alpha,\alpha_g) \\
&\frac{1}{\Delta^3}\mid_{u=\alpha_s+(1-v)\alpha_g} \\
&-8f_K m_b^2 m_K^4\int_0^1 dvv^2\int_0^1 d\alpha_g\int_0^{\alpha_g} d\beta\int_0^{1-\beta} d\alpha(A_\parallel+A_\perp)(1-\alpha-\beta,\alpha,\beta) \\
&\frac{1}{\Delta^3}\mid_{u=1-v\alpha_g}\ ,
\end{aligned}
\tag{20}
$$

where

$$
\begin{aligned}
\Delta &= m_b^2-(q+up)^2\,,\\
\Phi &= A_\parallel+A_\perp-V_\parallel-V_\perp\,.
\end{aligned} \tag{21}
$$

The light-cone distribution amplitudes of the K meson are defined by

$$
\begin{aligned}
\langle 0|\bar u(0)\gamma_\mu\gamma_5 s(x)|K(p)\rangle &= if_K p_\mu\int_0^1 due^{-iup\cdot x}\left\{\phi_K(u)+\frac{m_K^2x^2}{16}A(u)\right\}\\
&\quad +f_K m_K^2\frac{ix_\mu}{2p\cdot x}\int_0^1 due^{-iup\cdot x}B(u)\,,\\
\langle 0|\bar u(0)i\gamma_5 s(x)|K(p)\rangle &= \frac{f_K m_K^2}{m_s+m_u}\int_0^1 due^{-iup\cdot x}\varphi_p(u)\,,\\
\langle 0|\bar u(0)\sigma_{\mu\nu}\gamma_5 s(x)|K(p)\rangle &= i(p_\mu x_\nu-p_\nu x_\mu)\frac{f_K m_K^2}{6(m_s+m_u)}\int_0^1 due^{-iup\cdot x}\varphi_\sigma(u)\,,\\
\langle 0|\bar u(0)\sigma_{\alpha\beta}\gamma_5 g_s G_{\mu\nu}(vx)s(x)|K(p)\rangle &= f_{3K}\left\{(p_\mu p_\alpha g^\perp_{\nu\beta}-p_\nu p_\alpha g^\perp_{\mu\beta})-(p_\mu p_\beta g^\perp_{\nu\alpha}\right.\\
&\quad \left.-p_\nu p_\beta g^\perp_{\mu\alpha})\right\}\int\mathcal{D}\alpha_i\varphi_{3K}(\alpha_i)e^{-ip\cdot x(\alpha_s+v\alpha_g)}\,,\\
\langle 0|\bar u(0)\gamma_\mu\gamma_5 g_s G_{\alpha\beta}(vx)s(x)|K(p)\rangle &= f_K m_K^2 p_\mu\frac{p_\alpha x_\beta-p_\beta x_\alpha}{p\cdot x}\\
&\quad \int\mathcal{D}\alpha_i A_\parallel(\alpha_i)e^{-ip\cdot x(\alpha_s+v\alpha_g)}\\
&\quad +f_K m_K^2(p_\beta g_{\alpha\mu}-p_\alpha g_{\beta\mu})\\
&\quad \int\mathcal{D}\alpha_i A_\perp(\alpha_i)e^{-ip\cdot x(\alpha_s+v\alpha_g)}\,,\\
\langle 0|\bar u(0)\gamma_\mu g_s\tilde G_{\alpha\beta}(vx)s(x)|K(p)\rangle &= f_K m_K^2 p_\mu\frac{p_\alpha x_\beta-p_\beta x_\alpha}{p\cdot x}\\
&\quad \int\mathcal{D}\alpha_i V_\parallel(\alpha_i)e^{-ip\cdot x(\alpha_s+v\alpha_g)}\\
&\quad +f_K m_K^2(p_\beta g_{\alpha\mu}-p_\alpha g_{\beta\mu})\\
&\quad \int\mathcal{D}\alpha_i V_\perp(\alpha_i)e^{-ip\cdot x(\alpha_s+v\alpha_g)}\,,
\end{aligned} \tag{22}
$$

where the operator $\tilde G_{\alpha\beta}$ is the dual of the $G_{\alpha\beta}$, $\tilde G_{\alpha\beta}=\frac{1}{2}\epsilon_{\alpha\beta\mu\nu}G^{\mu\nu}$ and $\mathcal{D}\alpha_i$ is defined as $\mathcal{D}\alpha_i=d\alpha_1 d\alpha_2 d\alpha_3\delta(1-\alpha_1-\alpha_2-\alpha_3)$. The light-cone distribution amplitudes are

parameterized as

$$
\begin{aligned}
\phi_K(u,\mu) &= 6u(1-u)\left\{1+a_1C_1^{\frac{3}{2}}(2u-1)+a_2C_2^{\frac{3}{2}}(2u-1)\right\}\,, \\
\varphi_p(u,\mu) &= 1+\left\{30\eta_3-\frac{5}{2}\rho^2\right\}C_2^{\frac{1}{2}}(2u-1) \\
&\quad+\left\{-3\eta_3\omega_3-\frac{27}{20}\rho^2-\frac{81}{10}\rho^2a_2\right\}C_4^{\frac{1}{2}}(2u-1)\,, \\
\varphi_\sigma(u,\mu) &= 6u(1-u)\left\{1+\left[5\eta_3-\frac{1}{2}\eta_3\omega_3-\frac{7}{20}\rho^2-\frac{3}{5}\rho^2a_2\right]C_2^{\frac{3}{2}}(2u-1)\right\}\,, \\
\varphi_{3K}(\alpha_i,\mu) &= 360\alpha_u\alpha_s\alpha_g^2\left\{1+\lambda_3(\alpha_u-\alpha_s)+\omega_3\frac{1}{2}(7\alpha_g-3)\right\}\,, \\
V_\parallel(\alpha_i,\mu) &= 120\alpha_u\alpha_s\alpha_g\left(v_{00}+v_{10}(3\alpha_g-1)\right)\,, \\
A_\parallel(\alpha_i,\mu) &= 120\alpha_u\alpha_s\alpha_g a_{10}(\alpha_s-\alpha_u)\,, \\
V_\perp(\alpha_i,\mu) &= -30\alpha_g^2\left\{h_{00}(1-\alpha_g)+h_{01}\left[\alpha_g(1-\alpha_g)-6\alpha_u\alpha_s\right]\right. \\
&\quad\left.+h_{10}\left[\alpha_g(1-\alpha_g)-\frac{3}{2}\left(\alpha_u^2+\alpha_s^2\right)\right]\right\}\,, \\
A_\perp(\alpha_i,\mu) &= 30\alpha_g^2(\alpha_u-\alpha_s)\left\{h_{00}+h_{01}\alpha_g+\frac{1}{2}h_{10}(5\alpha_g-3)\right\}, \\
A(u,\mu) &= 6u(1-u)\left\{\frac{16}{15}+\frac{24}{35}a_2+20\eta_3+\frac{20}{9}\eta_4\right. \\
&\quad+\left[-\frac{1}{15}+\frac{1}{16}-\frac{7}{27}\eta_3\omega_3-\frac{10}{27}\eta_4\right]C_2^{\frac{3}{2}}(2u-1) \\
&\quad\left.+\left[-\frac{11}{210}a_2-\frac{4}{135}\eta_3\omega_3\right]C_4^{\frac{3}{2}}(2u-1)\right\}+\left\{-\frac{18}{5}a_2+21\eta_4\omega_4\right\} \\
&\quad\left\{2u^3(10-15u+6u^2)\log u+2\bar{u}^3(10-15\bar{u}+6\bar{u}^2)\log\bar{u}\right. \\
&\quad\left.+u\bar{u}(2+13u\bar{u})\right\}\,, \\
g_K(u,\mu) &= 1+g_2C_2^{\frac{1}{2}}(2u-1)+g_4C_4^{\frac{1}{2}}(2u-1)\,, \\
B(u,\mu) &= g_K(u,\mu)-\phi_K(u,\mu)\,,
\end{aligned}
\tag{23}
$$

where

$$
\begin{aligned}
h_{00} &= v_{00} = -\frac{\eta_4}{3}\,, \\
a_{10} &= \frac{21}{8}\eta_4\omega_4 - \frac{9}{20}a_2\,, \\
v_{10} &= \frac{21}{8}\eta_4\omega_4\,, \\
h_{01} &= \frac{7}{4}\eta_4\omega_4 - \frac{3}{20}a_2\,, \\
h_{10} &= \frac{7}{2}\eta_4\omega_4 + \frac{3}{20}a_2\,, \\
g_2 &= 1 + \frac{18}{7}a_2 + 60\eta_3 + \frac{20}{3}\eta_4\,, \\
g_4 &= -\frac{9}{28}a_2 - 6\eta_3\omega_3\,,
\end{aligned} \tag{24}
$$

here $C_2^{\frac{1}{2}}$, $C_4^{\frac{1}{2}}$ and $C_2^{\frac{3}{2}}$ are Gegenbauer polynomials, $\eta_3 = \frac{f_{3K}}{f_K}\frac{m_q+m_s}{m_K^2}$ and $\rho^2 = \frac{(m_u+m_s)^2}{m_K^2}$ [27, 28, 29, 30, 38, 39, 40, 41].

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